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4D gravity on a brane from bulk higher-curvature terms

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ABSTRACT: We study a gravity model where a tensionful codimension-one three-brane is embedded in a bulk with infinite transverse length. We find that 4D gravity is induced on the brane already at the classical level if we include higher-curvature (Gauss-Bonnet) terms in the bulk. Consistency conditions appear to require a negative brane tension as well as a negative coupling for the higher-curvature terms.

Contents

1. Introduction

In the context of brane world several mechanisms, that yield 4D gravity on a brane embedded in a larger space-time, have been studied. Here we would like to present some work done on a codimension-one setup where gravity is induced on the brane from bulk higher-curvature terms. In particular, we study a gravity model where a bulk five-dimensional Einstein-Hilbert (EH) action is ameliorated with the Gauss-Bonnet (GB) combination of curvature-squared terms. We also include a bulk cosmological constant and a 3-brane Σ with generically non-vanishing tension. Such a model has been thoroughly investigated over the past few years under many aspects, often viewed as a generalization [2, 3] of the RS setup [1]. In particular, in [3] it was shown that, upon adding GB to the lowest order action of [1], one obtains the Newtonian potential on the (positive tension) brane both in the IR and in the UV regime. Recently [4] it was found that negative tension solutions in the RS context may present tachyonic instabilities, in presence of a bulk GB term. Here, we study the aforementioned model under a different perspective. First of all we allow for solutions that have infinite invariant length in the transverse direction. In such a situation it is by now well known that 4D gravity can be induced on the brane through quantum effects [5, 6]. At the level of the low-energy effective action, if the localized matter is non-conformal, loops of matter fields with external gravity lead to a power series in derivatives of curvatures, that truncated at the second order amount to include on the brane an explicit EH term as well as a contribution to the brane tension [6]; it is usually referred to this model as the Dvali-Gabadadze-Porrati (DGP) model. By now several string theory realizations of the DGP model are known [7].

It was also pointed out that, upon including a bulk GB term, one can obtain 4D gravity on a tensionful codimension-2 brane in infinite transverse space, even without any explicit EH term on the brane [9].

In this note we study the occurrence of 4D gravity on a tensionful 3-brane in 5D bulk via the presence in the bulk of Gauss-Bonnet combination of curvatures. [1]

Hence the model reads [2]

$$S = \frac{1}{2\kappa^2} \int d^D x \ \sqrt{-G} \Big(R + \tilde{\xi} Z - \Lambda \Big) - \int_\Sigma d^{D-1}x \ \sqrt{-\hat{G}} f \qquad (1.1)$$

where

$$Z = R^2 - 4R_{MN}R^{MN} + R_{MNST}R^{MNST} \ , \qquad (1.2)$$

is the GB combination and

$$\hat{G}_{\mu\nu} = \delta_\mu^M \delta_\nu^N G_{MN} \Big|_\Sigma \ , \qquad (1.3)$$

is the pull-back of the metric G_{MN} on the brane. Although such a setup may present interesting cosmological features [2, 10, 11], here we limit our study to static solutions.

In order to stress the effect of the presence of the bulk GB term in our setup, we consider a toy-model where the Einstein-Hilbert term is absent, showing that one may get 4D gravity on a codimension-one brane from *pure* bulk quadratic terms.

The letter is organized as follows. In section 2 we describe the type of backgrounds upon which gravity has been studied. In section 3 some preliminary estimates concerning the parameters of the model are given. In section 4 we study linearized gravity on the static solutions described in section 2 and then we present some comments in the final section.

2. Static solutions

The equations of motion for the model (1.1) are given by

$$R_{MN} - \frac{1}{2}G_{MN} \Big(R + \tilde{\xi} Z - \Lambda \Big) + 2\tilde{\xi} Z_{MN} = -\frac{1}{2}G_{M\rho}G_{N\sigma}\hat{G}^{\rho\sigma}\frac{\sqrt{-\hat{G}}}{\sqrt{-G}}\tilde{f}\delta(z) \ , \quad (2.1)$$

where

$$Z_{MN} = RR_{MN} - 2R_{MS}R^S_{\ N} + R_{MRST}R_N^{\ RST} - 2R^{RS}R_{MRNS} \qquad (2.2)$$

and $\tilde{f} = 2\kappa^2 f$ (in the following we will also identify $\tilde{\xi} = 2\kappa^2\xi$).

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[1]In [8] a Gauss-Bonnet correction to the tensionless setup of [6] was considered.

[2]For notational convenience we work with a codimension-one brane on a space-time of unspecified dimension D, even though we will mostly have in mind the specific case $D = 5$.

We consider static warped solutions of the kind

$$ds^2 = e^{2A(z)}\eta_{MN}dx^M dx^N \tag{2.3}$$

where z parameterizes the transverse direction. The equations of motion are

$$(D-2)(A'' - A'^2)\left[1 - 2\tilde{\xi}(D-3)(D-4)A'^2 e^{-2A}\right] + \frac{1}{2}e^A\tilde{f}\delta(z) = 0 \tag{2.4}$$

$$(D-1)(D-2)\left[1 - \tilde{\xi}(D-3)(D-4)A'^2 e^{-2A}\right]A'^2 + \Lambda e^{2A} = 0 \tag{2.5}$$

which, as usual, admit bulk AdS solutions

$$A_\pm(z) = -\ln(\pm kz + 1) \qquad k > 0 \tag{2.6}$$

that, in turns, can be combined to give

$$A_{1,2}(z) = -\ln(\pm k|z| + 1) \tag{2.7}$$

in order to satisfy the jump condition included in (2.4). In fact, the equations of motion become

$$\tilde{f} = \pm 4(D-2)k\left[1 - 2\tilde{\xi}(D-3)(D-4)k^2\right] \tag{2.8}$$

$$k^2 - \tilde{\xi}(D-3)(D-4)k^4 + \frac{\Lambda}{(D-1)(D-2)} = 0 \tag{2.9}$$

and thus

$$k^2 = \frac{1}{2(D-3)(D-4)\tilde{\xi}}\left[1 + \epsilon\,\mathrm{sgn}(\xi)\sqrt{1 + 4\Lambda\tilde{\xi}\frac{(D-3)(D-4)}{(D-1)(D-2)}}\right] \qquad \epsilon = \pm 1 \tag{2.10}$$

are the generic bulk solutions for both $A_{1,2}$. The branch characterized by $\epsilon = -\mathrm{sgn}(\xi)$, in the limit $\xi \to 0$ reduces to the pure Einstein solution, and is referred to as the EH branch. The other branch ($\epsilon = +\mathrm{sgn}(\xi)$) is called GB branch and is not continuously connected to pure Einstein solutions.

The Ricci scalar on the solutions $A_{1,2}$ is given by

$$R = -D(D-1)k^2 \pm 4(D-1)k\delta(z) \tag{2.11}$$

The solution A_1 is therefore a GB deformation of the RS2 solution [1] and has finite invariant-length (compactification volume)

$$L \equiv \int_{-\infty}^{\infty} dz\ \sqrt{-G} = \frac{2}{(D-1)k} \tag{2.12}$$

although the range of the $z-$coordinate is infinite, as usual. On the other hand the lower solution, A_2, has infinite invariant-length

$$L \equiv \int_{-1/k}^{1/k} dz \, \sqrt{-G} = +\infty \, . \tag{2.13}$$

In general, the invariant length enters in the expression of the localized Planck mass as

$$M_P^{D-3} \sim L M^{D-2} \tag{2.14}$$

and thus the solution 2 cannot lead to localization of gravity. However the presence of the GB term turns out to play a special role here: it will in fact "induce" 4D gravity on the brane, along the lines of what described in [9]. However we will see that the present codimension-one case is quite peculiar in that positivity arguments appear to require a negative tension on the brane as well as a negative GB coupling. In the following we investigate some features of the aforementioned infinite-length solution and will argue that such apparently odd setup does not present evident inconsistencies.

3. The induced Planck mass and the brane tension

In order to show that the model (1.1) reproduces 4D gravity on the brane Σ also for the infinite-length solution, $A_2(z)$, we consider the equations of motion at the linearized level. It is not difficult to convince oneself that the GB combination induces a (D-1)-dimensional graviton propagator on the brane. The contribution of a term quadratic in curvatures to the linearized equations of motion can be schematically represented as

$$\xi R^{(0)} E_{\mu\nu}^{(1)} \sim \xi k \delta(z) \Box_4 h_{\mu\nu} \, , \tag{3.1}$$

where $E_{\mu\nu}^{(1)}$ is the linearized Einstein tensor and $R^{(0)}$ is the distributional part of the zero-th order Ricci scalar given in (2.11). One can thus recognize

$$M_P^{D-3} \sim -\xi k \tag{3.2}$$

and therefore the positivity of the induced Planck mass requires [3]

$$\xi < 0 \, . \tag{3.3}$$

From the jump condition we then see that [4]

$$f \sim -k[1 - 2\tilde{\xi}(D-3)(D-4)k^2] < 0 \tag{3.4}$$

[3]A model with negative Gauss-Bonnet coupling recently already appeared in a somewhat different context [11].

[4]In the tensionless limit equations (2.8-2.9) lead to the solitonic solution of [12]. Note that in such a case $\xi > 0$ and the positivity condition on M_P forces to choose the finite-length solution A_1.

and therefore both the tension and the GB coupling turn out to be negative in such a set up. Note that this result is valid only for the EH branch; in fact, in the GB branch, equation (2.10) would lead to negative k^2 as the coupling ξ is negative. On the other hand, in the EH branch, $k^2 > 0$ if the cosmological constant is negative.

4. Gravity on the brane

In order to have a more precise understanding of the mechanism that induces gravity in this setup let us consider the linearized equations of motion in presence of matter localized on the brane. To do this, let us study small fluctuations around the solution:

$$G_{MN} = \exp(2A)\left[\eta_{MN} + \widetilde{h}_{MN}\right] , \tag{4.1}$$

where for convenience reasons we have chosen to work with $\widetilde{h}_{MN}$ instead of metric fluctuations $h_{MN} = \exp(2A)\widetilde{h}_{MN}$. We will use the following notation

$$H_{\mu\nu} \equiv \widetilde{h}_{\mu\nu} , \quad A_\mu \equiv \widetilde{h}_{\mu D} , \quad \rho \equiv \widetilde{h}_{DD} . \tag{4.2}$$

for the component fields. The coupling between the localized matter and the graviton field reads

$$S_{\text{int}} = \frac{1}{2} \int_\Sigma d^{D-1}x \ T_{\mu\nu} H^{\mu\nu} . \tag{4.3}$$

We thus get the following set of equations

$$\left[1 - 2(D-3)(D-4)\tilde{\xi} A'^2 e^{-2A}\right] \left\{ \partial_\sigma \partial^\sigma H_{\mu\nu} + \partial_\mu \partial_\nu H - 2\partial_{(\mu}\partial^\sigma H_{\nu)\sigma} - \right.$$

$$\eta_{\mu\nu}\left(\partial_\sigma\partial^\sigma H - \partial^\sigma\partial^\rho H_{\sigma\rho}\right) + H''_{\mu\nu} - \eta_{\mu\nu}H'' + (D-2)A'\left(H'_{\mu\nu} - \eta_{\mu\nu}H'\right) -$$

$$2\left[\partial_{(\mu}A'_{\nu)} - \eta_{\mu\nu}\partial^\sigma A'_\sigma + (D-2)A'\left(\partial_{(\mu}A_{\nu)} - \eta_{\mu\nu}\partial^\sigma A_\sigma\right)\right] +$$

$$\left.\left[\partial_\mu\partial_\nu\rho - \eta_{\mu\nu}\partial_\sigma\partial^\sigma\rho + \eta_{\mu\nu}\left((D-2)A'\rho' + (D-1)(D-2)A'^2\rho\right)\right]\right\} -$$

$$4(D-4)\tilde{\xi}\left(A'' - A'^2\right)e^{-2A}\left\{\partial_\sigma\partial^\sigma H_{\mu\nu} + \partial_\mu\partial_\nu H - 2\partial_{(\mu}\partial^\sigma H_{\nu)\sigma} - \right.$$

$$\left.\eta_{\mu\nu}\left(\partial_\sigma\partial^\sigma H - \partial^\sigma\partial^\rho H_{\sigma\rho}\right) + (D-3)A'\left[H'_{\mu\nu} - 2\partial_{(\mu}A_{\nu)} - \eta_{\mu\nu}\left(H' - 2\partial^\sigma A_\sigma\right)\right]\right\} +$$

$$2(D-2)\left(A'' - A'^2\right)\left[1 - 4(D-3)(D-4)\tilde{\xi} A'^2 e^{-2A}\right]\eta_{\mu\nu}\rho$$

$$= -2\kappa^2\left(T_{\mu\nu} + \frac{1}{2}\eta_{\mu\nu}f\rho\right)\delta(z) , \tag{4.4}$$

$$\left[1 - 2(D-3)(D-4)\tilde{\xi} A'^2 e^{-2A}\right]\left[(\partial^\mu H_{\mu\nu} - \partial_\nu H)' - \partial^\mu F_{\mu\nu} + \right.$$

$$\left.(D-2)A'\partial_\nu\rho\right] = 0 , \tag{4.5}$$

$$\left[1 - 2(D-3)(D-4)\tilde{\xi}A'^2 e^{-2A}\right]\left[-(\partial^\mu \partial^\nu H_{\mu\nu} - \partial^\mu \partial_\mu H) + \right.$$

$$\left.(D-2)A'(H' - 2\partial^\sigma A_\sigma) - (D-1)(D-2)A'^2\rho\right] = 0 . \tag{4.6}$$

The *graviphoton* A_μ can be set to zero everywhere. The reason is twofold; on the one hand $A_\mu(z)$ is Z_2-odd and therefore it vanishes at $z = 0$; then using the unbroken diffeomorphisms A_μ can be gauged away completely [13]. On the other hand A_μ can couple to brane matter only through $\partial_\mu T^{\mu\nu}$ which vanishes for conserved brane matter.

Hence, setting such field to zero, eq. (4.5) reduces to

$$(\partial^\mu H_{\mu\nu} - \partial_\nu H)' + (D-2)A'\partial_\nu\rho = 0 \tag{4.7}$$

and can be solved to give

$$(H_{\mu\nu} - \eta_{\mu\nu}H)' + (D-2)A'\eta_{\mu\nu}\,\rho = \eta_{\mu\nu}F(z) . \tag{4.8}$$

It is not difficult to see that $F \equiv 0$ as the fluctuations are expected to vanish asymptotically. Therefore, using the background equation of motion (2.4), equation (4.4) can be cast in the form

$$\left[1 - 2(D-3)(D-4)\tilde{\xi}A'^2 e^{-2A}\right]\left\{\partial_\sigma\partial^\sigma H_{\mu\nu} + \partial_\mu\partial_\nu H - 2\partial_{(\mu}\partial^\sigma H_{\nu)\sigma} - \right.$$

$$\eta_{\mu\nu}\left(\partial_\sigma\partial^\sigma H - \partial^\sigma\partial^\rho H_{\sigma\rho}\right) + H''_{\mu\nu} - \eta_{\mu\nu}H'' - \left[\partial_\mu\partial_\nu\rho - \eta_{\mu\nu}\partial_\sigma\partial^\sigma\rho + \right.$$

$$\left.\left.\eta_{\mu\nu}(D-2)\left(A'\rho' + A'^2\rho\right)\right]\right\} - 4(D-4)\tilde{\xi}\left(A'' - A'^2\right)e^{-2A}\left\{\partial_\sigma\partial^\sigma H_{\mu\nu} + \right.$$

$$\left.\partial_\mu\partial_\nu H - 2\partial_{(\mu}\partial^\sigma H_{\nu)\sigma} - \eta_{\mu\nu}\left(\partial_\sigma\partial^\sigma H - \partial^\sigma\partial^\rho H_{\sigma\rho}\right)\right\} = -2\kappa^2\hat{T}_{\mu\nu}\delta(z) , \tag{4.9}$$

$$-(\partial^\mu\partial^\nu H_{\mu\nu} - \partial^\mu\partial_\mu H) + (D-2)A'H' - (D-1)(D-2)A'^2\rho = 0 , \tag{4.10}$$

where

$$\hat{T}_{\mu\nu} = T_{\mu\nu} - \frac{1}{2}\eta_{\mu\nu}f\rho . \tag{4.11}$$

In (4.9) the terms in the curly bracket multiplying

$$\left(A'' - A'^2\right)e^{-2A} = 2k\delta(z) \tag{4.12}$$

are nothing but the linearized Einstein tensor in (D-1)-dimensional flat space. Therefore it is easy to recognize

$$M_P^{D-3} = -8(D-4)k\xi \tag{4.13}$$

as the induced Planck mass on the brane. The leftover contribution in (4.9) is the bulk graviton propagator with an effective mass

$$2\kappa^2 M_{\text{eff}}^{D-2} = 1 - 2(D-3)(D-4)\tilde{\xi}A'^2 e^{-2A} = 1 - 2(D-3)(D-4)\tilde{\xi}k^2 \ . \quad (4.14)$$

In other words, the linearized equations of motion (4.9-4.10) for the spin-two fluctuations, are consistent with those associated to the model

$$S_* = M_{\text{eff}}^{D-2} \int d^D x \sqrt{-G} R + M_P^{D-3} \int_\Sigma d^{D-1}x \sqrt{-\hat{G}}\hat{R} + \frac{1}{2}\int_\Sigma d^{D-1}x T_{\mu\nu}H^{\mu\nu} \ . \quad (4.15)$$

At the linearized level the model (4.15) coincides with the model (1.1), albeit they differ beyond the linearized approximation. It is worth noting that, at the linearized level, the only effect of the GB combination in the bulk is the "renormalization" of the Planck mass. If fact, let us stress that one would obtain a canonical bulk propagator with positive nonvanishing bulk Planck mass even *without* the bulk Einstein-Hilbert term, that is starting from the purely higher-derivative model

$$S_{GB} = \int d^D x \ \sqrt{-G}\Big(\xi Z - V_B\Big) - \int_\Sigma d^{D-1}x \ \sqrt{-\hat{G}}f \ . \quad (4.16)$$

Such a Planck mass would be

$$M_{GB}^{D-2} = -2(D-3)(D-4)\xi k^2 \quad (4.17)$$

that is positive as the coupling ξ is negative in this model. In this limit, where the bulk EH term is absent

$$k = \left[\frac{V_B}{(D-1)\cdots(D-4)\xi}\right]^{\frac{1}{4}} \quad (4.18)$$

is the expression of the AdS scale as can be easily inferred from (2.9), and the brane Planck mass is again given by (4.13). This limit is unphysical as it is of course not obvious that higher curvature contributions would not spoil the results. However, we presented it here as we think it might be helpful to understand better how 4D gravity is induced on the brane in the setup we have considered.

Before ending this section we would like to point out some crucial differences with the results obtained recently by Deruelle and Sasaki [3]. The setup studied in [3] differs from the one studied here in that they considered a RS2 type of background, that is, 4D flat space with a noncompact warped finite-length extra dimension and a positive-tension 3-brane. In such a case two mechanisms are at work. On the one hand, there is a Randall-Sundrum type of localization that yields 4D gravity on the brane at large scales ($r >> r_1 \equiv 1/k$). On the other hand there is a Brane Induced Gravity (BIG) mechanism, given there by the bulk GB, that yields 4D gravity at small scales ($r << r_2 \sim \xi k/M^{D-2}$). Hence, in the BIG dominant regime

the crossover scale r_2 covers the other crossover scale r_1, and therefore perturbative gravity is 4D at *all* scales, as pointed out in [14]. In our model the invariant length of the transverse direction is infinite, and therefore we have no RS localization. We thus have a single BIG crossover scale

$$r_C \equiv \frac{-8(D-4)\xi k}{M^{D-2} - 2(D-3)(D-4)\xi k^2} \tag{4.19}$$

and gravity is four dimensional at distances much smaller than the previous scale.

5. Comments

The aim of this letter was to show that gravity can be "induced" on a 3-brane embedded in a 5D bulk with infinite transverse invariant length, via the presence of a bulk GB combination. In fact, for the only purpose of further underlining this mechanism, we have shown that gravity on the brane remains 4D in the ultraviolet regime even in a toy-model with purely GB bulk gravity. In [4] it was shown that negative-tension branes in the RS+GB context project-in a normalizable tachyon and project-out a (normalizable) zero-mode, thus leading to an instability of the setup. We claim that such arguments of instability do not apply in our case as gravity here is not localized or, in other words, there is no (normalizable) 4D graviton zero mode [5].

As to the negativity of the GB coupling, note first that $\xi < 0$ yields a positive contribution to the bulk kinetic term, so that there is no fear of bulk ghost fields. On the other hand, from string theory there is *a priori* no stringent constraint on the sign of the GB coupling. In fact, if on the one hand the sign of the Gauss-Bonnet combination in 10-dimensional heterotic string theory is positive [15], on the other hand in compactified theories such sign might depend on the details of the compactification. In fact, the Gauss-Bonnet combination in 5D can be, for instance, obtained from the compactification of M-theory on CY_3 [16]. In such a case the coupling constant for the GB term is moduli dependent and its sign is *a priori* not fixed.

Finally, a comment is in order regarding models with generic (non GB) combinations of quadratic curvature terms. Such models do not appear to admit solutions with delta-function type of discontinuities (thin sources), like the ones presented here. On the other hand it is plausible to conjecture that they might induce gravity on smooth types of branes. It is also plausible that such mechanism might persist in higher (larger than two) codimension setups. Curvature-squared terms were shown to be helpful to smooth out singularities in higher-codimension brane worlds with infinitely large transverse space [17]; they could also be beneficial to induce gravity in such setups. A thorough investigation in this direction is still missing.

Acknowledgments

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